

String/M Theory

C P Bachas

Laboratoire de Physique Théorique de l' Ecole Normale Supérieure
24 rue Lhomond, 75231 Paris Cedex 05, France

E-mails: bachas@physique.ens.fr, bachas@cpht.polytechnique.fr

Abstract

This is a brief review of the present status, of some recent developments and of the open challenges in string/M theory.[†]

1. Why string theory

The main claim of string theory is that it cures the singular short-distance behaviour of Einstein gravity, without giving up the fundamental laws of quantum mechanics. Although this is far from established, there are serious indications that the claim could indeed be true. In any case no other real alternatives to string theory are known at present. One can draw an analogy with spontaneously-broken gauge theories, which cured the singular short-distance behaviour of the four-Fermi interactions. The resolution of this theoretical inconsistency led to the construction of the Standard Model, and one hopes that the resolution of the problems associated with quantum gravity will, likewise, lead us to the fundamental microscopic theory that lies beyond.

A frequent objection is that while the Fermi scale was clearly accessible to accelerator experiments, quantum gravity need not manifest itself below $M_{\text{Planck}} \simeq 1,2 \times 10^{19} \text{ GeV}$, a scale surely out of direct experimental reach. String theory predicts, however, many types of new physics: supersymmetric partners, Kaluza-Klein and string excitations, enhanced gauge symmetry, and a larger gravitational sector. All of these are required for consistency, and the scales at which they should appear is an important open question to which I will return later in this talk. Even if one adopts the conservative hypothesis that non-gravitational physics much below M_{Planck} can be described by a minimal supersymmetric extension of the Standard Model (MSSM), it is increasingly hard to discard string theory as irrelevant. First, the MSSM is an effective theory with so many arbitrary parameters that it has little predictive power without extra assumptions. Second, physics at a superheavy

scale can have many indirect manifestations: we are used to associate neutrino masses, superlight axions, proton decay or primordial cosmology with the physics of conventional grand unification, and $M_{\text{GUT}} \simeq 2 \times 10^{16} \text{ GeV}$ is at most one order of magnitude below the fundamental string scale.

The search for a unified theory has been the main preoccupation of string theorists over the past fifteen years. Effective string excitations appear however also in many other contexts, and most notably in the low-energy limit of confining gauge theories like QCD. The early development of dual resonance models was in fact triggered by the observed Regge behavior of meson scattering, and the hope to find a controllable string approximation to QCD has been an important motivation in the subject ever since. The idea has been revived recently in a beautiful and unexpected way that I will discuss. Though it is too early to see where it will lead, it does illustrate how string theory has replaced quantum field theory as the new exciting theoretical frontier.

The plan of my talk is as follows: I will first briefly review the status of string theory up to the early 90's, and the rapid subsequent developments associated with duality symmetries, black holes and D-branes. I will then focus on three issues that have received attention over the past year: efforts to do away with supersymmetry, the AdS/CFT correspondence, and the question of string theory scales or of whether we live on a brane. This list is by no means exhaustive, it reflects partly my taste and partly my ignorance. There are some interesting topics I will not touch – cosmology or non-commutative theories to name a couple – and many more interesting papers I do not cite. Time limitation and ignorance is more frequently at cause here than taste.

[†] Plenary talk at the EPS HEP99 Conference, Tampere, Finland

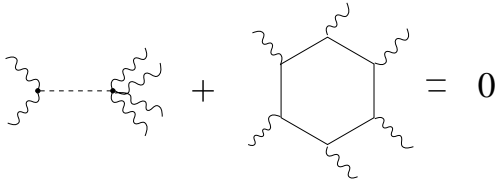


Figure 1. The Green-Schwarz anomaly-cancellation between a tree and a one-loop hexagon diagram.

2. Perturbative string theories

There are five known perturbative string theories, or rather consistent sets of Feynman rules (reviewed in many books on the subject [1]): the type IIA and IIB theories, the heterotic theories with gauge group $SO(32)$ or $E_8 \times E_8$, and the type I theory with gauge group $SO(32)$. Except for the type I theory whose quanta are unoriented closed and open strings, all the others have only oriented closed strings in their spectrum. They are all defined in ten space-time dimensions, and have N unbroken supersymmetries, where $N=2$ for the type II theories, and $N=1$ for the rest. One of their key ingredients is anomaly cancellation. Anomalies are quantum failures of classical symmetries, which are fatal when the symmetries in question are local. Anomaly cancellation requires for instance complete families of quarks and leptons in the Standard Model. The effective $N=1$ supergravity plus super Yang-Mills theories in ten dimensions were believed to be plagued by both gauge and gravitational anomalies [2]. It was the important discovery of Green and Schwarz [3] to show that string theory avoids the problem in a subtle way: the symmetry is violated classically, and restored by quantum corrections as illustrated schematically in figure 1.

Having found perturbatively consistent theories of gravity, the next step was to see if they could be made to fit our low-energy world. To do that one had to compactify six of the ten space-time dimensions, obtain a gauge group large enough to contain $SU(3)_c \times SU(2) \times U(1)$, and generate on the way chiral matter in which the quark and lepton families could fit. Now $N=1,2$ in ten dimensions corresponds to $N=4,8$ in four, and extended supersymmetry does not allow chirality in four dimensions (4d). In the process of compactification one should thus break either all, or all but one, of the original supersymmetries. For reasons that will become clear later this second alternative has been privileged, so that a lot of effort went into constructing ‘semirealistic’ $N=1$ compactifications to 4d. The two type II theories proved too restrictive to accommodate the Standard

Model [4], the type I theory was at the time technically more awkward to deal with, but for the heterotic theories many such compactifications were found and extensively analyzed in the mid eighties [5].

The main challenge of string theory still today is related to this embarrassment of riches. It is the question of *vacuum selection and stability*. In a nutshell, if a supersymmetric vacuum is a good first approximation to our world, there exist a zillion candidates for the job, and if not there are no stable candidates whatsoever. Part of the huge degeneracy of supersymmetric vacua is due to the existence of massless moduli. These are parameters of the compactification, which appear as 4d scalar fields with vanishing potential to all orders in the string coupling constant. Their expectation values are thus perturbatively undetermined – a nuisance since the masses, couplings and other properties of the effective theory change drastically as one moves around moduli space. Supersymmetry breaking, possibly triggered by non-perturbative gaugino condensation [6], could in principle lift this large degeneracy but it has proven hard to do reliable calculations in this context. Things are even worse for non-supersymmetric compactifications, where perturbation theory itself breaks down. Finding a stable non-supersymmetric vacuum in a theory containing gravity is of course a very ambitious task, since it amounts to solving among other things the problem of the cosmological constant.

3. Dualities and M theory

It is very likely that the solution to these long-standing problems will require a definition of string theory that goes beyond the Feynman-rules of the previous section. We are still lacking such a definition today, but progress towards this goal has been made. A key idea, which has opened a window into the non-perturbative structure of the theory, is the idea of duality symmetries. It can be illustrated with the Georgi-Glashow model, in which a $SO(3)$ gauge group breaks down to a $U(1)$ by the vacuum expectation value (v) of a higgs-triplet field.[†] At weak gauge coupling ($g \ll 1$) the spectrum of this theory contains, in addition to a neutral higgs, three other stable particles:

- a photon (γ)
- a charged gauge boson (W), and
- a magnetic monopole (M).

[†] This was an early candidate for a theory of the electroweak interactions, which failed to account for weak neutral currents.

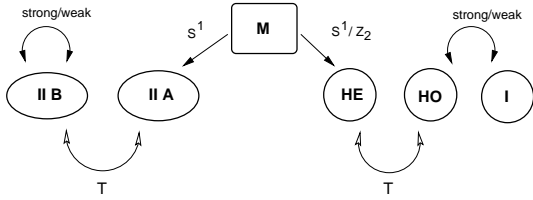


Figure 2. The web of string theory dualities.

The latter arises as a soliton (discovered by ‘t Hooft and Polyakov) and solitons are much heavier than elementary quanta,

$$m_W \sim gv \ll m_M \sim v/g. \quad (1)$$

The above inequality is valid at weak coupling, but if we allow $g \gg 1$ the roles of the electric and magnetic charges are naively exchanged. It is thus tempting to suppose that there exists a ‘dual’ perturbative formulation of the strongly-coupled theory, in which monopoles are the fundamental quanta (coupling with strength $\tilde{g} \sim 1/g$), and the W-bosons arise as solitonic excitations.

The problem with this kind of conjecture is that most of the time there is nothing one can calculate to test it. The mass formulae (1) for example receive in general both perturbative and non-perturbative corrections, so at strong coupling they need not be even qualitatively correct. There is however one notable exception: extended N=4 supersymmetry protects both the mass and the gauge coupling from quantum corrections, and the idea could in this case be put to the test with success [7]. As shown by Seiberg, Witten and others, electric/magnetic duality is in fact also a powerful tool in less supersymmetric, N=2 and N=1 gauge theories, but describing these celebrated results is outside the scope of the present talk (for reviews see for example [8]).

Dualities have their natural habitat in string theory [9] (for precursor ideas see [10], for reviews see [11]). They take many different forms: coupling constant inversion (S duality), inversion of the compactification volume in string units (T duality), or coupling/volume interchange (U duality). The five perturbative string theories of the previous section are related by such a web of dualities, as illustrated in figure 2. They are thus believed to be dual descriptions of the same physics, each best adapted in different regions of the moduli space of some fundamental underlying theory named M . Though we do not have a precise definition of M theory, one striking fact about it is that it exists in eleven flat space-time dimensions, where

its low-energy limit coincides with the maximal supergravity of Cremmer, Julia and Scherk [12]. Compactification of the eleventh dimension on a circle [9] or a line segment [13] leads to the type IIA or heterotic $E_8 \times E_8$ theories, at weak coupling when the circle or segment are very small.

Dualities have transformed our thinking about string theory in many ways. First they represent an extraordinary unification of principle: the theory of quantum gravity seems to be unique (there aren’t five different consistent theories), and even if vacuum selection and stability remains a puzzle, it is conceptually comforting that the issue is at least entirely dynamical. Second, they have uncovered previously inaccessible regions of moduli space, with new exciting phenomenological possibilities. Third, they have given us new tools for understanding the still mysterious physics of black holes. Finally, duality multiplets include besides fundamental strings a new class of excitations, whose role has been pervasive in the subject ever since: Dirichlet and other p-branes.

4. D-branes and black holes

A p-brane is an excitation with p-dimensional spatial extent: p=0 for particles, p=1 for strings, p=2 for membranes and so on (the list would stop there if we did not live in higher dimensions). Many conventional field theories have p-brane solitonic excitations. These are solutions of the classical non-linear field equations, localised in the transverse dimensions, and characterised by finite tension (mass/volume) and charge densities. The ‘t Hooft Polyakov monopole is such a soliton, other examples are the cosmic strings or domain walls of conventional grand unified models. It was the important insight of Polchinski [14] (for reviews see [15]) to recognize that type II string theories contain a particular class of such excitations that admit a very simple description as D(irichlet) branes. In a theory with only closed strings in the bulk, the (p+1)-dimensional worldvolume of a Dp-brane is a space-time defect on which fundamental open strings can end.

It turns out that this almost poetic definition determines the properties of the soliton precisely, with no need to ever solve any non-linear field equations explicitly.† Consider for example the interaction of two D-branes. To leading order

† The technical reason is that the classical string-field equations are equivalent to the requirement of (super)conformal invariance on the worldsheet, which is respected by the Dirichlet boundary conditions of string endpoints.

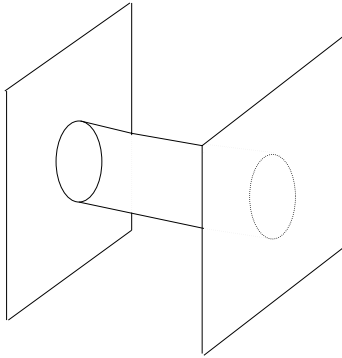


Figure 3. Leading interaction of two D-branes through the exchange of a closed string. The diagram has the dual interpretation as Casimir force – see text.

this is given by the exchange of a closed string (as in figure 3), whose massless modes are the graviton and its supersymmetric cohort. Only these contribute at large brane separation, but to evaluate the interaction we need to know how they couple to the solitons – we need to know in particular the D-brane charges and their tensions. The same diagram has however also a dual interpretation: it is a Casimir force due to the vacuum fluctuations of open strings, much like the Casimir force between two ordinary superconducting plates. This only depends on the spectrum of stretched open strings, so that the result can be evaluated readily [14][16] both for large brane separations and also at short distances, where the supergravity approximation is not valid.

The importance of D-branes derives from two basic facts: (a) they can ‘trap’ in their worldvolume non-abelian gauge fields, in addition to spin 1/2 and spin 0 excitations,[17] and (b) being solitons in a theory of gravity they can be assimilated in appropriate circumstances to large semiclassical black holes [18]. The first of these facts is illustrated in figure 4: the open strings living on a single D-brane reduce at low energies to an abelian gauge boson plus its supersymmetry partners. Putting n D-branes together gives rise to n^2 different types of open strings, whose low-energy excitations correspond now to a supersymmetric $U(n)$ gauge theory. Separating the branes breaks this gauge symmetry spontaneously to $U(1)^n$, while assembling different types of D-branes can reduce the number of unbroken supersymmetries. When n is large such collections of branes become heavier and heavier, and should eventually behave like large semiclassical black branes or black holes.

The super Yang-Mills/supergravity descriptions are a priori mutually exclusive, much like the two

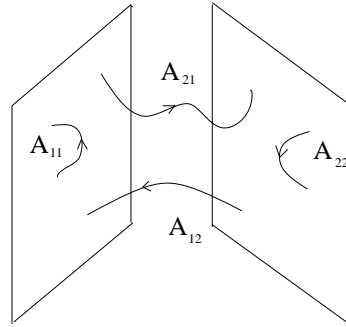


Figure 4. The four different types of open strings in the presence of two D-branes. Their low-energy limit corresponds to a $U(2)$ gauge theory.

dual descriptions of the Georgi-Glashow model. The argument for this goes roughly as follows: the mass of the would-be black hole in units of the fundamental string length is $M_{\text{BH}} \sim n/g_s$, where g_s is the string coupling constant and $\sim 1/g_s$ is the tension of a D-brane. The coupling of the effective super Yang Mills theory is $g_{\text{YM}}^2 \sim g_s$, while Newton’s constant is $G_N \sim g_s^2$ – this is because in the topological expansion a closed-string loop has double the weight of an open loop. Now for the supergravity description to be valid the Schwarzschild radius of the black hole must be much bigger than the string length,

$$M_{\text{BH}} G_N \sim n g_s \gg 1 \quad (\text{supergravity}) \quad (2)$$

(a more careful argument should take into account compactification radii). The Yang Mills description, on the other hand, is valid when the loop counting parameter is small

$$n g_{\text{YM}}^2 \sim n g_s \ll 1 \quad (\text{Yang – Mills}) \quad (3)$$

where n is the number of gauge bosons running in a loop. The two descriptions have thus non-overlapping ranges of validity.

Fortunately, supersymmetry can come once more to our rescue. For certain extremal black holes corresponding to supersymmetric configurations of branes, non-renormalization allows the extrapolation of some results from one limit to the other. This is true in particular for the entropy which can be calculated in the field theory at weak coupling, and then compared to the Beckenstein-Hawking area law [20] obtained by thermodynamic arguments on the black hole side. The agreement of the two results [18](for a detailed account see [19]) gave the first microscopic derivation of black hole entropy. This has boosted our confidence that string theory is a consistent microscopic theory of gravity even if, to be sure, the puzzles of black hole evaporation remain.

5. Is there life without supersymmetry?

Most of the discussion up to now has relied heavily on supersymmetry: duality checks, the microscopic black hole entropy calculation and vacuum stability, all depend crucially on it. Any sharp statement about non supersymmetric situations is thus a priori interesting, and some progress in this direction has been made over the past year, thanks mainly to the work of Sen (for a review and references see [21]). His idea was to consider non supersymmetric states (these are states in long or non-BPS multiplets of the supersymmetry algebra) which are protected from decaying by some discrete symmetry. The simplest example is provided by heterotic string states in spinor representations of $SO(32)$. The lightest such state must be stable even at strong coupling, and if duality is correct this state should exist in the dual type I description of the same theory. Sen identified this state as a controlled remnant in the annihilation of a D-brane anti D-brane pair, and succeeded to calculate its mass and other properties. Many more examples have since been found, and classified by the elegant mathematical theory called K-theory [22] (this is a topological classification, it has no claim to being a fundamental theory of nature like M !)

Can this help us to construct non-supersymmetric stable vacua? Despite some brave recent attempts in this direction [23] I don't think the answer is anywhere in sight yet. Let me recall what the situation was in older attempts to break supersymmetry in the heterotic string. The breaking scale could be (i) the heterotic string scale M_h [24], (ii) some Kaluza-Klein scale $1/r$, [25] or (iii) some dynamical scale $\sim M_h \exp(-1/g_h^2)$, if the breaking was induced by gaugino condensation [6] (here g_h is the heterotic string coupling). The fact that only the first two options are perturbative, motivated some early speculation that string theory should have TeV-size radii [26]. Weak coupling forces however, as I will explain in section 7, both M_h and $1/r$ to be near the Planck scale, so only the third possibility looks theoretically plausible. If one gives up low-energy supersymmetry altogether, vacuum instabilities at the string scale render perturbative calculations unreliable – even if the problem can sometimes be postponed [23]. In the gaugino condensation scenario things ‘look better’ but the important puzzles remain: why doesn't vacuum energy at the breaking scale drive the modulus g_h to zero, and why does the energy vanish at the minimum anyway?

The type I theory does seem to offer some novel possibilities. The vacua of this theory contain generally D-branes, as well as another type of

defect, or rather ‘end of the world’, called orientifold planes [27] (for reviews see for example [28]). These should fill all three non-compact dimensions of our universe, or else Poincare invariance would be lost, but they may be localized in the compact space. Supersymmetry can be broken on these branes and walls, while being to leading order preserved in the ten-dimensional bulk [29] (a similar phenomenon has been discussed in the Horava-Witten limit of the strongly-coupled heterotic theory [30]). Can we then use the stable non-BPS branes discussed above to construct a non-supersymmetric type I vacuum? The difficulty is that such branes carry in general energy and, as a result, gravitate. Since in a compact space ‘gravitational flux’ has nowhere to escape to, the total energy should vanish. This is however precisely our old friend the cosmological constant problem.

6. New hope for the QCD string ?

I have already mentioned in the introduction that the Regge behaviour of meson scattering was the earliest experimental manifestation of relativistic string theory, and the first motivation for developing dual resonance models. QCD has taken much of the steam out of this program, but has not lead to a quantitative control of strong interactions except at high energies where the theory is asymptotically free. Effective chiral lagrangians and lattice gauge theory provide useful approximations at lower energies, but the very property of confinement is a strong hint that a better string-theoretic description may exist. Past attempts to use the known string theories for QCD have failed in many respects:

- the critical dimension is 10 (or 26)
- there is a massless spin-2 particle
- what about parton behaviour?

In short, the known flat-space string theories have only a very remote resemblance to real QCD.

In a separate development ‘t Hooft pointed out [31] that Yang Mills theory retains its essential properties (confinement and asymptotic freedom) but simplifies considerably if one takes the number of colors n to be large. Keeping the combination

$$\lambda = ng_{YM}^2 \quad (4)$$

fixed in the limit, he showed that Feynman diagrams admit a topological expansion reminiscent of the loop expansion of string theory. Diagrams that can be drawn on a plane are weighted by n^2 , those that can be drawn on a torus by n^0 , and so on down the line. Of course $n = 3$ in real life, but this could still be a nice expansion with some luck!

Its leading term in which only the planar diagrams survive has, in particular, a ‘classical smell’ in it. The search for this classical limit (or ‘master field’) has been pursued over the years with little success. More recently, coming from different considerations, Polyakov proposed [32] that the QCD string must live in a curved geometry with more than four space-time dimensions.

These ideas have been revived and made very concrete in the last couple of years [33][34][35] (a complete review with an extensive list of references is [36]). Progress came from a closer study of the D-brane/black hole correspondence of the previous section, and was crystallized into a sharp conjecture by Maldacena [34]. His argument was to compare the following two different limits of the same system:

(i) in the D-brane description the limit where the fundamental string tension goes to infinity. What survives is a gauge field theory on the branes decoupled from the gravity in the bulk, schematically

$$(\text{brane YM}) \oplus (\text{bulk gravity}) \quad (5)$$

Gravity decouples because Newton’s constant has dimensions of inverse string tension to some positive power.

(ii) from the point of view of a distant observer, the limit in which his calorimeter can only measure vanishingly small energies. What our observer sees are very long-wavelength gravitons in the bulk, decoupled from *all* possible excitations near the horizon where they are redshifted to zero energy,

$$(\text{string near horizon}) \oplus (\text{bulk gravity}) \quad (6)$$

Equating (5) and (6) leads to the conjecture that the field theory on the brane has a dual description as a string theory in the near-horizon geometry of a black hole.

This correspondence is often called AdS/CFT because in the best-studied examples the field theory on the brane is conformal and the near horizon geometry is anti de Sitter. In the special case of n D3 branes the field theory is the $N=4$ supersymmetric version of $SU(n)$ Yang Mills, whose parameters are identified on the string side as follows

$$\lambda \sim (R/l_s)^4 \quad \text{and} \quad n \sim (R/l_P)^4 \quad (7)$$

Here R is the radius of the AdS geometry, l_s the string length scale and l_P the ten-dimensional Planck length. This identification renders clear the meaning of the ‘t Hooft expansion: $1/n$ corrections are string loop effects and as $n \rightarrow \infty$ the theory

becomes indeed classical – but this is a classical string theory in a curved higher-dimensional space-time! Now classical string theory is described by some conformal invariant 2d model – so you may think we are done – but the model corresponding to this particular background has proven hard to deal with [37]. Definite statements can thus be made only in the supergravity limit where one may ignore stringy effects. This amounts to taking $\lambda \rightarrow \infty$, which is strong coupling on the gauge theory side.

Now isn’t strong coupling all we are after in real QCD? Unfortunately not, the reason being very roughly as follows: in real QCD the coupling runs with energy, which would translate into varying curvatures as one moves radially in the near horizon geometry on the string side (for explicit realizations see [38][39][40]). Supergravity allows us to only control the region in which curvatures stay small, the rest has to be excised away. This corresponds to imposing an ultraviolet cutoff in the gauge theory, so what we really have is a cutoff theory at strong bare coupling. To go to the continuum limit we must scale the bare coupling to zero as we take the cutoff to infinity – but this drives us precisely to the region that supergravity does not control. We face therefore the same problem as lattice gauge theory, but transformed in a surprising way: to approach the continuum limit we need non-perturbative control over a two-dimensional σ -model. Whether this model will prove in the future tractable remains to be seen, but the ideas behind this surprising duality are most likely here to stay.

7. The case for traditional unification

Let me come finally to the question that is of more direct interest to the audience here: ‘what are the scales of string unification?’ or put more provocatively, ‘will we see strings and extra dimensions at the LHC?’ (this is also discussed in Peskin’s talk [41]) The conventional (and conservative) hypothesis is that the string, compactification and Planck scales lie all to within two or three orders of magnitude from each other, and are hence far beyond direct experimental reach. The non-gravitational physics at lower energies is thus described by a 4d supersymmetric quantum field theory (SQFT), which must at least include in it the MSSM. This conventional hypothesis is supported by the following three solid facts : (i) Softly broken SQFTs *can* indeed be extrapolated consistently to near-Planckian energies without destabilizing the electroweak scale ; (ii) the hypothesis is (almost) automatic in the weakly-coupled heterotic string theory, and (iii) the minimal (or ‘desert’) string-unification assumption

is in remarkable agreement with some of the measured low-energy parameters of our world.

The first fact is the well-known reason for introducing supersymmetry in particle physics in the first place. Softly-broken supersymmetry solves the technical aspect of the gauge-hierarchy problem by preventing radiative corrections from driving M_{weak} to the ultraviolet cutoff. I will return to this point later.

The second argument is less known to non-string theorists and goes as follows: in the weakly-coupled heterotic theory both the graviton and the gauge bosons are massless modes of a closed string. They thus live in ten-dimensional spacetime and interact at tree-level through the sphere diagram. The effective four-dimensional Yang-Mills and Einstein actions, obtained by dimensional reduction, therefore read

$$\mathcal{L}_{\text{gauge}} \sim \frac{(rM_h)^6}{g_h^2} \text{tr} F^2 \quad (8)$$

and

$$\mathcal{L}_{\text{grav}} \sim \frac{r^6 M_h^8}{g_h^2} \mathcal{R} , \quad (9)$$

where M_h and g_h are the heterotic string scale and string coupling constant, while r is the typical compactification radius. From the coefficients of these actions one can read the four-dimensional gauge coupling and Planck mass,

$$g_{\text{YM}} \sim \frac{g_h}{(rM_h)^3} , \quad (10)$$

and

$$M_h \sim g_{\text{YM}} M_{\text{Planck}} . \quad (11)$$

Factors of 2's and π 's in these relations are irrelevant to our arguments and have been dropped. Suppose now that g_{YM} is of order one, or 1/10 it does not really matter. Then the universal relation (11) tells us that the string scale is tied automatically to the Planck mass [42]. Furthermore in view of relation (10) we must have rM_h also of order one, or else the string theory will be strongly coupled. [†] Thus all scales are tied up together and there is little leeway for abandoning the traditional hypothesis within the weakly-coupled heterotic theory.

Let me come finally to the third argument in favour of the traditional hypothesis. It is a well-known observation [43, 44] (for more recent discussions see [45]) that if we extrapolate the three measured low-energy gauge couplings with the β -functions of the MSSM, they meet at a scale $M_U \simeq$

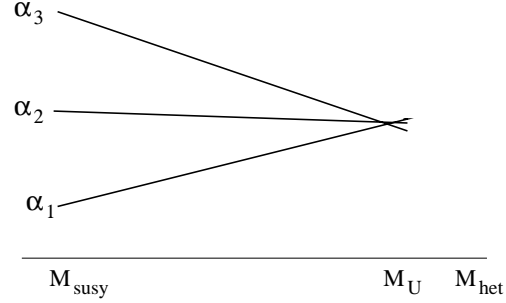


Figure 5. Schematic drawing of gauge coupling unification in the MSSM. On a logarithmic scale the discrepancy between the ‘experimental’ and the string-theory unification scale is a few percent.

2×10^{16} GeV. This is consistent with the one-loop formulae

$$\alpha_i^{-1}(\mu) = \alpha_U^{-1} + b_i \log \frac{\mu}{M_U} + \Delta_i \quad (12)$$

where b_i are the β -function coefficients, and α_U is the fine structure constant at M_U . The threshold corrections Δ_i parametrize our ignorance of the details of the theory at the unification scale, and of the details of supersymmetry breaking. Assuming that the Δ_i are negligible, and treating M_U and α_U as input parameters, we have one prediction which is verified by LEP data at the level of a few percent. Note that it is the existence of the energy desert which renders the above prediction meaningful and robust. Threshold effects make a few-percent correction to differences of couplings only because the logarithm in equations (12) is very large.

Gauge coupling unification is an indirect hint for the existence of the heavy scale M_U , and also for low-energy supersymmetry, since it works much better with the MSSM β -functions than with those of the SM. It is also a hint in favour of the traditional string unification, because M_U is remarkably close – though not exactly equal – to M_h . Since this latter is related to the experimentally known Planck mass and to α_U , it is in principle calculable. A careful calculation [46] gives $M_h \simeq 5 \times 10^{17}$ GeV. On the appropriate logarithmic scale this differs from the ‘experimental value’ M_U by a few percent, a rather remarkable a priori agreement[†] illustrated schematically in figure 5. After all there was no a priori reason why the extrapolated low-energy couplings should not have met, if at all, at say 10^{35} GeV !

[†] T-duality tells us that the radii cannot be smaller than the string length.

[¶] For a detailed discussion of the small discrepancy see [47]

8. Do we live on a brane?

I have already mentioned that dualities have uncovered some previously inaccessible regions of the moduli space of M-theory. The description of the vacuum in these regions involves various branes that can trap non-abelian gauge fields in their worldvolume: D-branes, and their dual heterotic fivebranes [48] and Horava-Witten walls [13]. Such vacua could realize the old idea [49] according to which we might be living on a brane, which traps in it the Standard Model fields.[†] Since gauge fields and gravity live now in different spaces, the universal relation tying up string, compactification and Planck scales is generally lost.

Let me illustrate this in type I theory, whose compactifications involve as we saw collections of D-branes and orientifolds. The graviton (a closed-string state) lives in the ten-dimensional bulk, and interacts to leading order through the sphere diagram. The open-string gauge bosons on the other hand are localized on the worldvolume of some collection of branes (I will call them our brane world) and interact via the disk diagram, which is higher order in the topological expansion. The Yang Mills and Einstein actions in four dimensions read

$$\mathcal{L}_{\text{gauge}} \sim \frac{(rM_I)^{6-n}}{g_I} \text{tr} F^2 \quad (13)$$

and

$$\mathcal{L}_{\text{grav}} \sim \frac{r^{6-n} \tilde{r}^n M_I^8}{g_I^2} \mathcal{R} , \quad (14)$$

where r is the typical radius of the n compact dimensions transverse to our brane world, \tilde{r} the typical radius of the remaining $(6-n)$ compact longitudinal dimensions, M_I the type-I string scale and g_I the string coupling constant. By T-dualities we can ensure that both r and \tilde{r} are greater than or equal to the fundamental string length.

Unlike the case of the heterotic string there is no universal relation between M_{Planck}^2 (the coefficient of the Einstein term), M_I and g_{YM}^2 (the inverse coefficient of the Yang Mills term). The radius r of the transverse space, in particular, enters only in the gravitational action, so by taking it to be large we can make the gravitational interaction weaker and weaker, while keeping M_I , g_I and g_{YM} fixed. This is explained intuitively by figure 6. Thus type I string theory is much more flexible (and less predictive) than its heterotic counterpart. This added flexibility can be used to ‘remove’ the order-of-magnitude discrepancy between the apparent

[†] A more extreme idea is that gravity itself could be trapped on a brane [50].

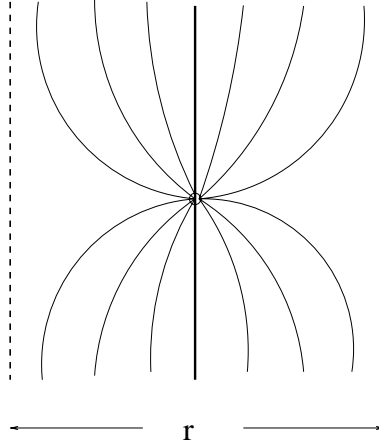


Figure 6. The spreading of flux in the transverse compact dimensions that could be responsible for the weakness of gravitational forces in a brane world.

unification and string scales [51]. If one decides however to ignore gauge coupling unification (efforts to reconcile matters are made in [55]) as accidental there is no reason why not to lower M_I further, to an intermediate [54] or even to the TeV scale [52][53]. Keeping for instance g_I and $(\tilde{r}M_I)$ fixed and of order one, leads to the condition

$$r^n \sim M_{\text{Planck}}^2 / M_I^{2+n} . \quad (15)$$

A TeV string scale would then require from $n = 2$ millimetric to $n = 6$ fermi-size dimensions transverse to our brane world.

What focused attention [52] to the brane-world idea was (i) the realization that it cannot be generally ruled out by existing data, even in the most extreme case of ‘TeV-ish’ string scale and millimeter-size transverse space, and (ii) the hope that this extreme case could be a new solution to the problem of the gauge hierarchy. Let me start by explaining the first point. Gravity is hard to test at submillimeter distances because of the large background of residual electromagnetic interactions. The ratio for instance of the Van der Waals to Newtonian force between two hydrogen atoms a distance d apart is

$$\frac{F_{\text{VdW}}}{F_{\text{grav}}} \sim \left(\frac{1\text{mm}}{d} \right)^5 . \quad (16)$$

At $d = 10\mu\text{m}$ Newton’s force is thus ten orders of magnitude weaker than Van der Waals. As a result the present-day data [56] allows practically any modification of Newton’s law, as long as it is of comparable strength at, and screened beyond, the

millimeter range. This has been appreciated in the past in the context of gravitational axions and light string moduli [57], and a similar bound holds for Kaluza-Klein modes.

Besides mesoscopic gravity experiments, one may worry about bounds from precision observables of the Standard Model, and from various exotic processes. Precision tests of the SM and compositeness bounds cannot however rule out in a model-independent way any new physics above the TeV-ish scale. Bounds for instance from LEP data on four-fermion operators, or bounds on dimension-five operators contributing to the $g - 2$ of the electron/muon are safe, as long as the characteristic scale of the new physics is a few TeV [52]. Proton decay and other exotic processes could of course rule out large classes of low-scale models, but natural suppression mechanisms do exist. One type of model-independent exotic process is graviton emission in the bulk, which could be seen as missing-energy events in collider experiments. The process is however suppressed by the four-dimensional Newton constant at low energies, and only becomes appreciable (as one should expect) near string scale where quantum gravity effects are strong [58] (see also Peskin's talk [41]).

Let me come finally to question of the gauge hierarchy. Although the brane world with a TeV string scale does not really solve the problem, it does transform it in an interesting way. Instead of asking why is M_{Planck} so much bigger than the SM scale, one now asks why is the transverse size r so much larger than the string length? In the particular case of an effectively two-dimensional transverse space, there is furthermore logarithmic sensitivity on r , which is (superficially) the natural setup for obtaining a hierarchy of scales [59]. It is, however, very hard to know whether any of this constitutes progress or not, without fully addressing the issue of vacuum stability.

9. Outlook

The titles of three out of the last four sections had a questionmark in them, and it is often said that 'if you see a question asked in a title the answer is NO'. Well, maybe so, but I am ready to bet that there *is* life without supersymmetry – we are after all the living proofs of it. As for the QCD string, even if this new attempt to master it fails, the ideas that underly it have already had a profound impact in the field and are most likely here to stay. What about the brane world and TeV strings? The recent developments have sunk in the point that we really know very little about

gravity. They have also brought string theorists and phenomenologists closer together, which is a marvelous thing. They have not yet, however, solved any theoretical problem, and abandoning the traditional hypothesis does make gauge coupling unification – our one solid empirical hint about physics beyond the SM – look accidental. As theorists we should thus keep an open-minded but cool head towards the possibility, which surely deserves further scrutiny. Ultimately it is vacuum selection and stability, or the LHC, that will of course decide!

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